

Hyperon-Nucleon Interaction in a Quark Model *

MAKOTO OKA[†]*Department of Physics, Tokyo Institute of Technology
Meguro, Tokyo 152, Japan*

ABSTRACT

A realistic hyperon (Y)-nucleon (N) interaction based on the quark model and the one-boson-exchange potential is constructed. The Nijmegen potential model D with the $SU(3)$ flavor symmetry is modified with a quark exchange interaction at the short-distance, which replaces the short-range repulsive core in the original model. The flavor-spin dependences of the short-range repulsion are qualitatively different from the original hard-core potential. We also study a two-body weak decay, $\Lambda N \rightarrow NN$, in the quark model. An effective weak interaction, where one-loop QCD corrections are explicitly taken into account, is employed. Differences from the conventional meson-exchange processes are discussed.

1. Introduction

In this lecture, I discuss two subjects: First, I present our recent attempt to constructing a realistic hyperon-nucleon interactions in the quark model. This part is based on the work done with Kenichiro Ogawa and Sachiko Takeuchi[1]. The second part is devoted for the study of nonmesonic weak decays of Λ in the direct quark processes, which is done in collaboration with Takashi Inoue and Sachiko Takeuchi[2].

2. $Y - N$ interaction

The short-distance repulsions between baryons seem universal for most two-baryon interactions. It is, for instance, known from the study of hypernuclei that the hyperon-nucleon interactions contain a short-range repulsion similar to the nuclear force. Why are the baryon-baryon interactions mostly repulsive at short distances? The simple quark model provides us with two possible answers to this question. Namely, there are two mechanisms in the quark model which produce short-distance repulsion[3,4].

*A lecture given at the International School Seminar on *Hadrons and Nuclei from QCD*, Tsuruga-Vladivostok-Sapporo, August-September, 1993.

[†]e-mail: oka@phys.titech.ac.jp

The first (I) is due to the Pauli exclusion principle among the valence quarks. This can be most easily demonstrated in an extreme case, $\Delta^{++}(S_z = \frac{3}{2})$ vs. $\Delta^{++}(S_z = \frac{1}{2})$, where all the quarks are UP in the flavor and spin \uparrow except for one quark in the second Δ . In a simple harmonic oscillator quark model, if two baryons stick to each other with relative 0s harmonic oscillator state, one has to excite at least two of the UP \uparrow quarks to a higher single particle orbit in order to satisfy the Pauli exclusion principle. Mathematically, it means that the antisymmetrized product of two Δ^{++} 's in relative 0s state vanishes:

$$\mathcal{A}|\Delta^{++}(\frac{3}{2})\Delta^{++}(\frac{1}{2})\chi_{0s}(R)\rangle = |u^6, S_z = 2, (0s)^6\rangle = 0 \quad (1)$$

where \mathcal{A} is the quark antisymmetrization operator for all the six quarks and χ_{0s} denotes the relative $\Delta\Delta$ wave function in the harmonic oscillator 0s state. Nonvanishing states require at least $2\hbar\omega$ excitation, which corresponds to the relative 1s state of two ground state. Namely, the relative 0s state of $\Delta^{++} - \Delta^{++}$ is forbidden and all the allowed $\chi(R)$ has to be orthogonal to χ_{0s} . Thus the relative wave function always has a node at $R = R_c$, which is nearly independent of the energy and the effective potential has a repulsive core of radius R_c .

One can interpret the above feature in a slightly different way. By multiplying $\langle\Delta^{++}\Delta^{++}\delta(R-S)|$ from the left of eq(1), one obtains

$$\begin{aligned} \langle\Delta^{++}\Delta^{++}\delta(R-S)| & |\mathcal{A}|\Delta^{++}\Delta^{++}\chi_{0s}(R)\rangle \\ &= \int \langle\Delta^{++}\Delta^{++}\delta(R-S)|\mathcal{A}|\Delta^{++}\Delta^{++}\delta(R-S')\rangle \chi_{0s}(S') dS' \\ &= \int N(S, S') \chi_{0s}(S') dS' = e \chi_{0s}(S) \end{aligned} \quad (2)$$

where $N(S, S')$ is the normalization integral kernel of the resonating group method and e is the eigenvalue of N associated with the eigenstate χ_{0s} . The forbidden state yields $e = 0$, while one obtains $e = 1$ if no antisymmetrization is considered. In general, the eigenvalue e gives a good indication of the “forbiddenness” of the two-baryon system. Namely, if $e < 1$, the channel has a “partially forbidden” state and the baryonic potential has a repulsion at short distances (or actually $R = 0$).

When the same argument is applied to the hyperon-nucleon systems, one finds that two $N\Sigma$ channels, $N\Sigma$ ($S = 0, I = \frac{1}{2}$) and $N\Sigma$ ($S = 1, I = \frac{3}{2}$), have small eigenvalues, $e = 1/9$ and $2/9$ respectively. They thus have an almost forbidden state. This indicates a strong repulsion in the S wave $N\Sigma$ interactions in those channels[5].

The second mechanism (II) for the short range repulsion is driven by the hyperfine interaction among quarks. The success of the quark model description of the meson-baryon spectrum owes largely to the spin-spin interaction

$$V_{\text{CMI}} = -\frac{\alpha_s}{4} \sum_{i < j} \frac{2\pi}{3m_i m_j} (\lambda_i \cdot \lambda_j) (\vec{\sigma}_i \cdot \vec{\sigma}_j) \delta(\vec{r}_{ij}) \quad (3)$$

which is considered to come from the magnetic part of a gluon exchange between quarks. The importance of this interaction in the baryon spectrum is manifested, for instance, in $N - \Delta$, and $\Lambda - \Sigma$ mass differences, and the negative neutron mean charge square radius.

The importance of the hyperfine interaction in the short-range NN interaction has been pointed out in the quark cluster model calculation[3,5]. One finds that the spin-spin interaction (3) produces a short-range repulsion not only for NN but also for other baryon-baryon interactions, such as $N\Lambda$ and $N\Sigma$. Such calculations also indicate that the Pauli exclusion principle (mechanism I) gives in general a stronger short-range repulsion than the hyperfine interaction (II).

3. Quark cluster model with the Nijmegen meson exchange potential

We concentrate on the hyperon-nucleon (YN) interaction here and present a realistic YN interaction model, which incorporates the quark exchange interaction at short distances and the meson exchange potential at larger distances[1]. The antisymmetrization of six valence quarks with the one-gluon exchange interaction leads to a strong repulsion, whose range is determined by the size of the baryon, and beyond its range the conventional meson-exchange processes take over and yield medium-long range attraction which binds nucleons together into nuclei.

We follow the $SU(3)$ symmetry for the meson-baryon couplings. Indeed, the YN potential models, such as the Nijmegen models[6] and Jülich models[7], are based on the $SU(3)$ symmetry. In this study, we employ the meson-exchange part of the Nijmegen potential model D and instead of using the hard cores in the original model, superpose it with the quark exchange interaction at the short distance.

The quark exchange interactions can be calculated in the quark cluster model (QCM) approach[3]. We consider a valence quark model with a hamiltonian,

$$H = K + V_{\text{CONF}} + V_{\text{OGE}} \quad (4)$$

where K is the nonrelativistic quark kinetic energy term, V_{CONF} stands for a quark confinement potential and V_{OGE} is the Fermi-Breit potential for the one gluon exchange. We employ the resonating group method (RGM) wave function for the six-quark system, given by

$$\Phi_{BB'}(1 \sim 6) = \mathcal{A}[\phi_B(1 \sim 3) \phi_{B'}(4 \sim 6) \chi(R)] \quad (5)$$

and the integral equation, called the RGM equation, with kernels H (Hamiltonian) and N (Normalization):

$$\int [H(R, R') - E N(R, R')] \chi(R') dR' = 0 \quad (6)$$

are solved. Nonlocality of the RGM equation comes from the antisymmetrization of the quarks. So far, we have not included effects of the instanton induced interactions in this model[8,9].

We introduce to the QCM equation (6) the meson exchange potential, which is borrowed from the Nijmegen model D in this study. This can be done by adding an integral kernel for the meson exchange potential, given by

$$V(R, R') \equiv \int dR'' N^{1/2}(R, R'') V_f(R'') N^{1/2}(R'', R') \quad (7)$$

where V_f is the Nijmegen meson exchange potential with the appropriate form factor. The form factor is chosen so as to be consistent with the quark wave function of the baryon,

$$V_f(R) \equiv \int \rho(x; R/2) V_N(x - y) \rho(y; -R/2) dx dy \quad (8)$$

where V_N is the Nijmegen potential without the repulsive core and the quark density of the baryon centered at $R/2$ is denoted by $\rho(x; R/2)$. In the QCM calculation, we employ the Gaussian wave function for the quark for simplicity, and thus the corresponding form factor is given also by a Gaussian.

We have five parameters in the model: the light quark mass m_q , the ratio of the light and strange quark masses m_q/m_s , the strength of confinement a , the strength of the one-gluon exchange potential α_s , and the size parameter b for the Gaussian wave function of quarks in the baryon. In order to make the calculation consistent in kinematics, we choose m_q to be one-third of the average octet baryon mass, *i.e.*, 383.7 MeV. The ratio of the light/strange quark masses is fixed to 0.6, which gives the $\Lambda - \Sigma$ mass difference. The gluon coupling constant is chosen so as to reproduce the $N - \Delta$ mass difference, and we also choose the confinement a so that the baryon state is stable against the breathing mode excitation, *i.e.*, $\partial E_B / \partial b = 0$. The remaining parameter b is sensitive to the NN interaction, because it determines the size of the form factor and also the range of the quark exchange interaction. Therefore we leave this as a free parameter and use the NN scattering data to choose the best value for b . The QCM calculation with the Nijmegen D meson exchange potential can fit the NN 1S_0 scattering phase shift well for $b = 0.56$ fm. Then the other parameters are determined: $a = 20.8$ MeV/fm, $\alpha_s = 1.85$.

We calculate the scattering S matrices for various YN systems in this model and find that the qualitative predictions given above are confirmed in the present model. In Table 1, we summarize the properties of the short-distance YN interactions. The lowest eigenvalue of the normalization kernel for each channel, given in the Table, distinguishes the first (I) and the second (II) mechanisms for the short-range repulsion. One sees that the Pauli exclusion principle gives a stronger repulsion for the $N\Sigma$ ($S = 0, I = \frac{1}{2}$) and $N\Sigma$ ($S = 1, I = \frac{3}{2}$) channels, while the other channels show a mild repulsion which is generally softer than the original Nijmegen model D. The repulsions in these two $N\Sigma$ channels are as strong as the Nijmegen model F, which is known to provide not enough

Table 1: Eigenvalues of the normalization kernel and the effective core radii for various S-wave YN systems. The “type” indicates the origin of the repulsion, either from the first (I) or the second (II) mechanisms. The effective core radii are obtained from the scattering phase shifts in the present model.

BB'	(J,I)	e	type	effective core radius
$N\Lambda$	$(0, \frac{1}{2})$	1	II	0.40 fm
$N\Sigma$	$(0, \frac{1}{2})$	$\frac{1}{9}$	I	0.68 fm
$N\Lambda$	$(1, \frac{1}{2})$	1	II	0.34 fm
$N\Sigma$	$(1, \frac{1}{2})$	1	II	0.30 fm
$N\Sigma$	$(0, \frac{3}{2})$	$\frac{10}{9}$	II	0.48 fm
$N\Sigma$	$(1, \frac{3}{2})$	$\frac{2}{9}$	I	0.67 fm

binding for Σ to make a bound Σ hypernuclei. Details of the model and the results will be published elsewhere[1].

4. Two-body Weak Decay of Λ

The hyperon Λ decays weakly into a nucleon and a pion in the free space. It, however, is suppressed in the nuclear medium by the Pauli blocking on the final nucleon state, whose momentum is less than 100 MeV/c for the Λ decay at rest. Indeed, in heavy hypernuclei, the decay is predominantly the nonmesonic one, that is, $\Lambda N \rightarrow NN$. If we assume that the initial Λ and the nucleon are at rest, then the final relative momentum of NN is about 420 MeV/c and thus is well above the Fermi momentum.

Theoretical study of the $\Lambda N \rightarrow NN$ decay has traditionally employed the meson (π , ρ , etc.) exchange mechanism, where one of the meson-baryon vertices involves the weak transition $s \rightarrow d$ [10]. Contributions from the direct quark-quark weak interaction, $us \rightarrow ud$ or $ds \rightarrow dd$, have not been taken into account. However, such direct quark-quark processes may play significant roles, as the relative NN momentum in the final state is not small.

Recent analyses of experimental data of decays of hypernuclei have revealed some difficulties in the meson-exchange picture. For instance, the so-called $n - p$ ratio, i.e., the ratio R_{np} of $\Lambda n \rightarrow nn$ v.s. $\Lambda p \rightarrow np$ decay in the nucleus, is predicted very small, $R_{np} \simeq 0.1$ in the meson-exchange picture. This is due to the strong contribution of the tensor force, which is preferred at the large momentum transfer. The tensor force selects the $S = 1$, $I = 0$ pn final state and therefore R_{np} becomes small. The experimental

data, however, seem not to agree with the prediction, i.e., $R_{np} \simeq 1$ in decays of light hypernuclei. We argue that the direct quark process, which does not follow the $I = 0$ selection rule, may enhance the $n - p$ ratio.

The mesonic weak decays of hyperons have been tested for the $|\Delta I| = \frac{1}{2}$ rule and are known to satisfy the rule to about 5% error. The same rule for the nonmesonic weak processes, like $\Lambda N \rightarrow NN$, is not confirmed yet. Indeed, an analysis of the decay of the $A = 3$ and 4 hypernuclei claims that the $|\Delta I| = \frac{1}{2}$ rule is not satisfied[11]. It is therefore urgent to clarify the mechanism of the $|\Delta I| = \frac{1}{2}$ rule in the free hyperon decays and to study whether the same mechanism restricts the nonmesonic decays to $|\Delta I| = \frac{1}{2}$ as well.

In the study of the meson-exchange processes, the $|\Delta I| = \frac{1}{2}$ rule is assumed from the beginning, implemented in the $\Lambda \rightarrow N\pi$ vertex. We instead employ the effective quark-quark weak hamiltonian, which contains both the $|\Delta I| = \frac{1}{2}$ and $|\Delta I| = \frac{3}{2}$ components. Although the $|\Delta I| = \frac{3}{2}$ part has a small overall coefficient, we will see that the matrix elements for the $\Lambda N \rightarrow NN$ decay may not be small compared to the $|\Delta I| = \frac{1}{2}$ component.

The effective weak hamiltonian describing $\Delta S = \pm 1$ processes, given by several authors[12,13], is

$$H_{eff}^{\Delta S=1} (Q^2 \sim \mu^2) = -\frac{G_f}{\sqrt{2}} \sum_{r=1, r \neq 4}^6 K_r O_r \quad (9)$$

where the four-quark operators, O_k ($k = 1, 2, 3, 5$ and 6) are defined by[12]

$$O_1 = (\bar{d}_\alpha s_\alpha)_{V-A} (\bar{u}_\beta u_\beta)_{V-A} - (\bar{u}_\alpha s_\alpha)_{V-A} (\bar{d}_\beta u_\beta)_{V-A} \quad (10)$$

$$O_2 = (\bar{d}_\alpha s_\alpha)_{V-A} (\bar{u}_\beta u_\beta)_{V-A} + (\bar{u}_\alpha s_\alpha)_{V-A} (\bar{d}_\beta u_\beta)_{V-A} \\ + 2(\bar{d}_\alpha s_\alpha)_{V-A} (\bar{d}_\beta d_\beta)_{V-A} + 2(\bar{d}_\alpha s_\alpha)_{V-A} (\bar{s}_\beta s_\beta)_{V-A} \quad (11)$$

$$O_3 = 2(\bar{d}_\alpha s_\alpha)_{V-A} (\bar{u}_\beta u_\beta)_{V-A} + 2(\bar{u}_\alpha s_\alpha)_{V-A} (\bar{d}_\beta u_\beta)_{V-A} \\ - (\bar{d}_\alpha s_\alpha)_{V-A} (\bar{d}_\beta d_\beta)_{V-A} - (\bar{d}_\alpha s_\alpha)_{V-A} (\bar{s}_\beta s_\beta)_{V-A} \\ = \frac{1}{3} O_2 + O_3(\Delta I = \frac{3}{2}) \quad (12)$$

$$O_5 = (\bar{d}_\alpha s_\alpha)_{V-A} (\bar{u}_\beta u_\beta + \bar{d}_\beta d_\beta + \bar{s}_\beta s_\beta)_{V+A} \quad (13)$$

$$O_6 = (\bar{d}_\alpha s_\beta)_{V-A} (\bar{u}_\beta u_\alpha + \bar{d}_\beta d_\alpha + \bar{s}_\beta s_\alpha)_{V+A} \quad (14)$$

The coefficients (Table 2) for the above six four-quark operators are calculated by using the renormalization group technique within the one-loop QCD corrections included[13]. The most prominent feature of this effective hamiltonian is that the QCD correction enhances the O_1 component while the other terms are suppressed. This is the main mechanism for the $\Delta I = 1/2$ enhancement.

This effective hamiltonian has been used for the calculations of the nonleptonic decay of strange mesons and baryons.[12,13,14] It is found that although the $\Delta I = 1/2$ enhancement is significant in those decays, agreement to experiment is not always achieved

Table 2: Strengths of the weak effective four-fermi vertexes

K_1	K_2	K_3	K_5	K_6
-0.265	0.010	0.026	0.003	-0.020

quantitatively. Some suggest that the decay amplitudes are sensitive to the meson and baryon wave functions.

We will adopt the above effective hamiltonian, but also consider the case where the $\Delta I = 3/2$ component of the operator O_3 is omitted and compare the results in order to see the effect of the $\Delta I = 3/2$ contribution.

In calculating the decay amplitude for $\Lambda N \rightarrow NN$, we employ the quark cluster picture for the two baryon systems. In the present calculation[2], we choose the most simple wave functions for the initial and the final states. First, we assume that the baryon consists of three valence quarks, whose orbital wave function is a harmonic oscillator eigenstate.

$$\phi(1, 2, 3)^{\text{orb}} = \left(\frac{1}{2\pi b^2}\right)^{\frac{3}{4}} \left(\frac{2}{3\pi b^2}\right)^{\frac{3}{4}} \exp\left\{-\frac{1}{4b^2}\vec{\xi}_{12}^2\right\} \exp\left\{-\frac{1}{3b^2}\vec{\xi}_{12-3}^2\right\} \quad (15)$$

The six quark wave function is given by

$$|\Lambda N\rangle = \mathcal{A}^6 |\phi(1, 2, 3)\phi(4, 5, 6)\chi_0(\vec{R})\rangle \quad (16)$$

$$|NN\rangle = \mathcal{A}^6 |\phi(1, 2, 3)\phi(4, 5, 6)\chi(\vec{R}')\rangle \quad (17)$$

where \mathcal{A}^6 is the antisymmetrization operator for six quarks, and ϕ is the internal wave function of the baryon. The flavor-spin wave function of the baryon is taken to be purely the SU(6) wave functions, which is known to be a good approximation.

In this article, we only consider the simplest case, in which the initial Λ and N are on top of each other and therefore the orbital part of the initial wave function is that of the $(0s)^6$ configuration in the harmonic oscillator shell model. Similarly, the final state is assumed simply a plane wave of two nucleon clusters.

$$\chi_0(\vec{R}) = \left(\frac{3}{2\pi b^2}\right)^{\frac{3}{4}} \exp\left\{-\frac{3}{4b^2}\vec{R}^2\right\} \quad (18)$$

$$\chi(\vec{R}') = \left(\frac{1}{2\pi}\right)^{\frac{3}{2}} \exp\{i\vec{k} \cdot \vec{R}'\} \quad (19)$$

Although these choices of the wave functions are not realistic, they will clarify the qualitative difference between the meson-exchange and the direct-quark processes, which is the purpose of this preliminary study. A full-range calculation using the realistic two-baryon wave functions is underway.

Table 3: Possible initial and final quantum numbers for the initial $L = 0$ transition and the calculated transition invariant matrix elements in $10^{-9} \text{ MeV}^{-1/2}$.

channel	isospin	spin-orbital		full	$\Delta I = \frac{3}{2}$ omitted
1	$p\Lambda \rightarrow pn$	$^1S_0 \rightarrow ^1S_0$	a_p	2.68	5.52
2		$^1S_0 \rightarrow ^3P_0$	b_p	2.07	0.48
3		$^3S_1 \rightarrow ^3S_1$	c_p	6.67	6.67
4		$^3S_1 \rightarrow ^1P_1$	e_p	-0.39	-0.39
5		$^3S_1 \rightarrow ^3P_1$	f_p	-1.31	-1.22
6	$n\Lambda \rightarrow nn$	$^1S_0 \rightarrow ^1S_0$	a_n	9.80	7.80
7		$^1S_0 \rightarrow ^3P_0$	b_n	-0.45	0.68
8		$^3S_1 \rightarrow ^3P_1$	f_n	-1.66	-1.72

Because we employ the nonrelativistic valence quark picture for the wave functions, the effective hamiltonian is also approximated by adopting the Breit-Fermi nonrelativistic expansion up to $1/c$.

If we restrict our initial state to $L = 0$, there exist eight possible combinations, given in Table 3, of L , S , and J for the initial and final states. We note that the $I = 1$ final states are allowed both for $(\Lambda n \rightarrow nn)$ and $(\Lambda p \rightarrow pn)$, while the $I = 0$ states are not possible for $(\Lambda n \rightarrow nn)$. Thus we have 5 $(\Lambda p \rightarrow pn)$ and 3 $(\Lambda n \rightarrow nn)$ matrix elements, which are labeled from a through f in Table 3, according to a widely used notation[15]. The results of the calculation are also given in Table 3. Eight amplitudes give all the information for the $\Lambda N \rightarrow NN$ weak decay in the present calculation. One sees that the parity conserving matrix elements (a and c) are dominant both in $I = 0$ and $I = 1$ channels. The last column of Table 3 shows the results after omitting the $|\Delta I| = \frac{3}{2}$ components of the O_3 operator. We find a significant contribution of the $|\Delta I| = \frac{3}{2}$ matrix elements.

In Table 4, we summarize the calculated decay rates with the initial spin averaged and the final states summed up. The $n - p$ ratio, $R_{np} \equiv \Gamma_n/\Gamma_p$, the ratio of the parity violating (PV) v.s. the parity conserving (PC) contributions, η , and the decay asymmetry parameter, a_1 , for the $\Lambda p \rightarrow pn$ and the $\Lambda n \rightarrow nn$ decays are also given in Table 4. We find that the $n - p$ ratio in the present calculation is much larger than that obtained in the meson-exchange calculation. It is very encouraging. Although the present calculation assumes a very naive wave function for the initial and final states, one sees at least the direct-quark process has the right direction to improve the meson exchange result, which

Table 4: Calculated observables

	full	$\Delta I = \frac{3}{2}$ omitted
$\Gamma_p (10^8 \text{sec}^{-1})$	2.82	3.16
$\Gamma_n (10^8 \text{sec}^{-1})$	1.96	1.31
R_{np}	0.70	0.42
η_p	0.070	0.031
η_n	0.088	0.15
$a_1(p)$	-0.28	-0.24
$a_1(n)$	0	0

is too small to account for the experimental value.

The decay asymmetry parameter describes the angular distribution of the outgoing two nucleons in the rest frame,

$$W(\theta) = 1 + a_1 \mathcal{P}_\Lambda P_1(\cos \theta) \quad (20)$$

where \mathcal{P}_Λ is the polarization of the lambda particle in the nucleus. Recent experiment done at KEK indicates a large a_1 for light hypernuclei. The data is consistent with $a_1 \simeq -1.0 \pm 0.4$. Our calculation yields the correct sign, but the magnitude seems too small.

5. Conclusion and Discussion

We present a quark model analyses of the hyperon-nucleon systems both for the strong and weak interaction processes. The realistic strong YN interactions, which are nonlocal due to the quark antisymmetrization effects, are proposed using the quark cluster model approach with the Nijmegen model D meson exchange potential. The main difference between the original Nijmegen model and our interaction arises in the spin-isospin dependence of the YN short range interactions. Especially, ΣN with $S = 0$, $I = 1/2$ and $S = 1$ and $I = 3/2$ have strong repulsion at the short distance in the quark model and may make the bound Σ hypernuclei unplausible.

We also present a quark model calculation of the direct-quark processes for the weak $\Lambda N \rightarrow NN$ decay, which can be observed exclusively in decays of hypernuclei. Assuming simple initial and final wave functions, we find that the calculated decay rates are comparable to the meson exchange contributions in magnitudes and show qualitatively

distinctive properties. This is encouraging because the direct-quark processes may resolve the discrepancies between experiment and the calculated results in the meson-exchange mechanism. The further study of the weak process is underway[2].

References

- [1] K. Ogawa, S. Takeuchi and M. Oka, to be published.
- [2] T. Inoue, S. Takeuchi and M. Oka, to be published.
- [3] M. Oka and K. Yazaki, Phys. Lett. **B90**(1980)41; Prog. Theor. Phys. **66**(1981)556; *ibid.* **66**(1981)572; in *Quarks and Nuclei*, ed. by W. Weise (World Scientific, 1985); K. Shimizu, Rep. Prog. Phys. **52**(1989)1; S. Takeuchi, K. Shimizu and K. Yazaki, Nucl. Phys. **A504**(1989)777.
- [4] M. Oka, in *Proceedings of 10th International Symposium on High Energy Spin Physics*, Nagoya, Nov., 1992.
- [5] M. Oka, K. Shimizu and K. Yazaki, Phys. Lett. **B130**(1983)365; Nucl. Phys. **A464**(1987)700; M. Oka, Phys. Rev. **D38**(1988)298.
- [6] M.M. Nagels, T.A. Rijken and J.J. de Swart, Phys. Rev. **D12**(1975)744; *ibid.* **15**(1977)2547; *ibid.* **20**(1979)1633; *ibid.* **17**(1978)768; P.M.M. Maessen, Th.A. Rijken and J.J. de Swart, Phys. Rev. **C40**(1989)2226.
- [7] B. Hilzenkamp, K. Holinde and J. Speth, Nucl. Phys. **A500**(1989)485; A. Reuber, K. Holinde and J. Speth, Jülich preprint, KFA-IKP(TH)-1993-06.
- [8] M. Oka and S. Takeuchi, Phys. Rev. Lett. **63**(1989)1780.
- [9] M. Oka and S. Takeuchi, Nucl. Phys. **A524**(1991)649; S. Takeuchi and M. Oka, Phys. Rev. Lett. **66**(1991)1271.
- [10] H. Bando, T. Motoba and J. Zofka, Int. Jour. Mod. Phys. **A5**(1990)4021, and references therein; A. Ramos, C. Bennhold, E. van Meijgaard, and B.K. Jennings, Nucl. Phys. **A547**(1992)103c.
- [11] R.A. Schumacher, Nucl. Phys. **A547**(1992)143c.
- [12] A.I. Vainshtein, V.I. Zakharov and M.A. Shifman, Sov. Phys. JETP, **45**(1977)670; L.B. Okun, *Leptons and Quarks* (North Hollandm The Netherland, 1982).
- [13] E.A. Paschos, T. Schneider, and Y.L. Wu, Nucl. Phys. **B332**(1990)285.
- [14] Y. Abe *et al.*, Prog. Theor. Phys. **64**(1980)1363.
- [15] M.M. Block and R.H. Dalitz, Phys. Rev. Lett. **11**(1963)96.